

# Three-dimensional MHD simulations of radiatively cooling, pulsed jets

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## ABSTRACT

We here investigate, by means of fully three-dimensional (3-D) Smoothed Particle Magnetohydrodynamic (SPMHD) numerical simulations, the effects of different initial magnetic field configurations on the evolution of overdense, radiatively cooling, pulsed jets, using different initial magnetic field topologies: (i) a longitudinal, (ii) a helical geometry permeating both the jet and the ambient medium, and (iii) a purely toroidal geometry permeating the jet only. We explore the effects of different pulsational periods, as well as different values of the magnetic field strength ( $\beta \simeq 0.1 - \infty$ , or  $B \simeq 260 \mu\text{G}-0$ ). The presence of a helical or a toroidal field tends to affect more the global characteristics of the fluid than a longitudinal field. However, the relative differences which have been previously detected in 2-D simulations involving distinct magnetic field configurations are diminished in the 3-D flows. While the presence of toroidal magnetic components can modify the morphology close to the jet head inhibiting its fragmentation in the early evolution of the jet, as previously reported in the literature, the impact of the pulsed-induced internal knots causes the appearance of a clumpy, complex morphology at the jet head (as required by the observations of Herbig-Haro jets) even in the MHD jet models with helical or toroidal configurations. The detailed structure and emission properties of the internal working surfaces can be also significantly altered by the presence of magnetic fields. The increase of the magnetic field strength (decrease of  $\beta$ ) improves the jet collimation, and amplifies the density (by factors up to 1.4, and 4) and the  $\text{H}\alpha$  intensity (by factors up to 4, and 5) behind the knots of jets with helical field and  $\beta \simeq 1 - 0.1$  (respectively), relative to a non-magnetic jet. As a consequence, the corresponding  $I_{[\text{S II}]} / I_{\text{H}\alpha}$  ratio (which is frequently used to determine the excitation level of HH objects) can be largely decreased in the MHD models with toroidal components relative to non-magnetic calculations. We also find that the helical mode of the Kelvin-Helmholtz instability can be triggered in MHD models with helical magnetic fields, causing some wiggling of the jet axis. No evidence for the formation of the nose cones that are commonly detected in 2-D jet simulations with initial toroidal magnetic fields,

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is found in the 3-D flows, nor even in the  $\beta \simeq 0.1$  case. The implications of our results for Herbig-Haro jets are briefly discussed.

*Subject headings:* ISM: jets and outflows – ISM: Herbig-Haro objects – MHD – star: formation

## 1. Introduction

The bright Herbig-Haro (HH) objects associated with low-mass young stellar objects, are radiating shock fronts immersed in collimated optical jets that may extend from a few 1000 AU (e.g., the microjets in the Orion Nebula; Bally, O’Dell & McCaughrean 2000), to very large parsec-length scales, like the *giant* HH jets (Heathcote et al. 1998; Bally & Devine 1994; see also Reipurth 1999 for a catalogue of such objects).

Currently, the most accepted model for the production of the discrete HH knots is by time variability in the driving sources of the jets. Their bow shock morphology, high spatial velocity, and symmetric pattern with respect to the star (e.g., Reipurth 1989a, 1989b; Bürke, Mundt & Ray 1988; Zinnecker, McCaughrean & Rayner, 1997, 1998) all indicate that the HH shocks generally arise through the steepening of velocity fluctuations in the underlying, supersonic outflow. Strong support for this conjecture has been given by theoretical studies which have confirmed that traveling shocks created in this manner reproduce the essential properties of the observed knots (e.g., Raga et al. 1990; Raga & Kofman 1992; Kofman & Raga 1992; Stone & Norman 1993; de Gouveia Dal Pino & Benz 1994; Völker et al. 1999; de Gouveia Dal Pino 2001).<sup>2</sup>

Magnetic fields seem to play a fundamental role in several phases of the jet launching and propagation processes. They seem to be relevant to: (i) operate locally in an accretion disk, in order to regulate the angular momentum process (Balbus & Hawley 1998) and the accretion rate; (ii) launch outflows through magneto-centrifugal forces from the disk itself (e.g., Blandford & Payne 1982; Kudoh, Matumoto & Shibata 1998; Königl & Pudritz 2000), or from the disk-star boundary region (the so-called X-point; e.g., Shu et al. 2000); and (iii) collimate the outflow once launched (e.g., Ouyed, Pudritz & Stone 1997). On the larger scales, since the outflowing material does not seem to be highly resistive (e.g., Frank et al. 1999), the magnetic field may also play a relevant role on its dynamics and evolution. Previously, Morse et al. (1992) have inferred an upper limit for the magnetic field in the ambient medium upstream the HH34 bow shock of about  $10 - 20 \times 10^{-6}$

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<sup>2</sup>We note that although pinch-modes of the Kelvin-Helmholtz (K-H) were originally suggested as the source of the knotty pattern in HH jets, it has been found that these modes have their strength highly reduced in the presence of radiative cooling (e.g., Blondin, Fryxell and Königl 1990; de Gouveia Dal Pino & Benz 1993). Besides, they typically lead to a too fast periodicity to match the observations (e.g., Stone, Xu, and Hardee 1997). Nevertheless, nonaxisymmetric K-H modes appear to be responsible for the gentle wandering or wiggling observed in many jets (e.g., de Gouveia Dal Pino, Birkinshaw, & Benz 1996; Stone, Xu, & Hardee 1997).

G. More recently, Ray et al. (1997), based on polarization measurements, have detected magnetic field strengths of the order 1 G at a distance of few tens of AU from T Tau S. Using magnetic flux conservation, this latter result would give magnetic fields of the order of  $\sim 10^{-3}$  G for a toroidal field configuration, and  $\sim 10^{-6}$  G for a longitudinal field at distances  $\sim 0.1$  pc. For typical jet parameters, this would imply a plasma parameter  $\beta = p_{j,gas}/(B^2/8\pi) \simeq 10^{-3}$  to  $10^3$  at those distances. Although Ray et al. have concluded that the high values measured for B at  $\sim 10$  AU probably come from regions of strong magnetic field amplification (such as behind shocks), the figures above suggest that magnetic fields must play some role in the overall jet dynamics and propagation even on the outer scales where local amplification (particularly behind the leading and the internal bow shocks) is expected to occur (e.g., Cerqueira, de Gouveia Dal Pino & Herant 1997; hereafter, Paper I).

Lately, there has been increasing focusing on MHD studies of the propagation of overdense, radiative cooling jets in an attempt to look for possible signatures of magnetic fields on the large scales of the HH outflows (see, e.g., de Gouveia Dal Pino & Cerqueira 2001 for a review). Frank et al. (1997, 1998), for example, in axisymmetric studies of steady-state jets with toroidal magnetic fields (and  $\beta = 8\pi p_{th}/B^2 < 1$ ), have reported the formation of pinch modes driven by magnetic tension, and the appearance of a magnetically confined nose cone structure of compressed gas between the bow shock and the Mach disk at the jet head.

In previous investigations (de Gouveia Dal Pino & Cerqueira 1996; Paper I; Cerqueira & de Gouveia Dal Pino 1999, hereafter Paper II), we have employed a three-dimensional Smoothed Particle Hydrodynamic (SPH) code that has been implemented to compute the effects of magnetic fields to investigate steady-state jets. Assuming different initial magnetic field configurations (in approximate equipartition with the gas) permeating both the jet and the ambient medium, we have found that there is an increase in the jet collimation in comparison to purely hydrodynamic calculations, which is caused by the amplification and reorientation of the magnetic fields behind the shocks (particularly in the presence of a helical field). Further, we have found that the presence of a helical magnetic field may inhibit the fragmentation of the dense shell formed by cooling and Rayleigh-Taylor instability at the head of the jet (Paper I), while a longitudinal field retains the fragmentation. This result could be an indication that the latter configuration could dominate near the jet head as the observed jets have a clumpy structure there. We have also evidenced from the simulations the development of the MHD K-H modes, but the weakness of the pinch modes makes it doubtful that they could play an important role in the production of the bright knots (Paper II). Where there is overlap, these results have been confirmed by recent axisymmetric calculations (e.g., Gardiner et al. 2000; Stone & Hardee 2000; O’Sullivan & Ray 2000). In particular, Gardiner et al. (2000), assuming longitudinal magnetic fields (with  $\beta = 0.1 - 10^7$ ), have found that the global characteristics of the flow are not strongly affected by the **B**-field strength, and a predominantly axial field tends to increase the collimation and order, and inhibit instabilities in the flow, but the presence of internal pulses increases the likelihood of magnetic reconnection.

The 2-D calculations of steady-state and pulsed MHD jets of Stone & Hardee (2000; hereafter,

SH) propagating into unmagnetized ambient media have also shown that the magnetic fields affect the fragmentation of the shell at the jet head, and a strong toroidal field peaking near the jet surface develops a nose cone which should be unstable in three-dimensions. They have also found that the radial hoop stresses due to the toroidal field confine the shocked jet material in the internal pulses resulting in higher densities in the pulses which are strongly peaked towards the jet axis in comparison to purely hydrodynamic calculations. A similar line of numerical studies has also been conducted by O’Sullivan & Ray (2000; hereafter OR). Most of their results essentially confirm previous hydrodynamic (e.g., Blondin, Fryxell, & Königl 1990; de Gouveia Dal Pino & Benz 1993, hereafter GB93; Stone & Norman 1994) and the MHD results above. However, unlike these works, they have assumed jets with initially lower densities with respect to the ambient medium (corresponding to a density ratio  $\eta = 1$ ), and which are highly over-pressurized with respect to the ambient medium. This causes the development of more complex cocoons surrounding the beams and the formation of crossing shocks which help to refocus the beam and the internal pulses. In particular, they find that the hoop stresses associated with toroidal fields can cause the *disruption* of the internal knots, even for  $\beta \simeq 1$ .

In a recent study, we have presented the first results of 3-D MHD simulations of the early evolution of pulsed jets (Cerqueira & de Gouveia Dal Pino 2001, hereafter, Paper III). In the present paper, we attempt to extend these previous investigations by exploring in three-dimensions the role played by three different initial magnetic field configurations on the evolution and emission structure of radiatively cooling, pulsed jets, considering a more extensive range of parameters. As before (Paper I, II and III), together with a baseline non-magnetic calculation, we employ a modified version of our 3-D SPH magnetized code and consider a longitudinal and a helical field topology (in approximate equipartition with the gas), both permeating the jet and the ambient medium. Also, in order to make a closer comparison with previous studies, a third geometry involving a purely toroidal field permeating the jet is considered.

The paper is planned as follows. In §2, we will briefly outline the numerical technique and setup. In §3, we present the results of the simulations for pulsed jets; in §4, we compare the results with previous 2-D MHD calculation and discuss some observational implications of them, and present our final remarks and conclusions.

## 2. Numerical method

We solve the equations that describe the evolution of a fluid using the Smoothed Particle Hydrodynamic (SPH) technique (e.g.; Benz 1990; Monaghan 1999). SPH is a Lagrangean scheme that makes use of particles to follow fluid parcels, and the evaluation of a given physical quantity is given by an interpolation scheme, using for this a kernel function ( $W$ ). The kernel has an adaptative smoothing length ( $h$ ) value, which defines the space volume of the SPH particles. Like any other physical quantities, magnetic fields can also be handled with SPH. In order to take into account the presence of the magnetic field, we have modified a pure HD, SPH code. In Paper II, we give the

details of such modification, as well as the results of some tests in the MHD limit. The application of the SPH technique to pure hydrodynamic jet simulations can be seen in several papers (e.g., GB93; de Gouveia Dal Pino & Benz 1994; Chernin et al. 1994; de Gouveia Dal Pino & Birkinshaw 1996; de Gouveia Dal Pino, Birkinshaw, & Benz 1996; de Gouveia Dal Pino 1999, 2001), so that our code has been continuously tested in both HD and MHD regimes. Such an MHD, SPH based code is often referred as “SPMHD” (e.g., Stellingwerf & Peterkin 1990; Meglicki 1994).

The solved system of MHD equations, in its ideal approximation, is given by:

$$\frac{d\rho}{dt} = -\rho \nabla \cdot \mathbf{v} \quad (1a)$$

$$\frac{d\mathbf{v}}{dt} = -\frac{\nabla p}{\rho} + \frac{1}{4\pi\rho}(\nabla \times \mathbf{B}) \times \mathbf{B} \quad (1b)$$

$$\frac{du}{dt} = -\frac{p}{\rho}(\nabla \cdot \mathbf{v}) - \mathcal{L} \quad (1c)$$

$$\frac{d\mathbf{B}}{dt} = -\mathbf{B}(\nabla \cdot \mathbf{v}) + (\mathbf{B} \cdot \nabla)\mathbf{v} \quad (1d)$$

where  $\rho$  is the density;  $\mathbf{B}$  is the magnetic field;  $u$  is the specific internal energy and  $\mathcal{L}$  is the radiative cooling rate. As in our previous works (see Paper II and references therein), we have adopted the cooling function given by Katz (1989) for a gas of cosmic abundances cooling from  $T \simeq 10^6\text{K}$  to  $\simeq 10^4\text{K}$ . An ideal equation of state is used to close the above system of equations:

$$p = (\gamma - 1)\rho u \quad (1e)$$

with  $\gamma = 5/3$ .

## 2.1. Initial and boundary conditions

Our computational domain is a rectangular box with dimensions  $-30R_j \leq x \leq 30R_j$ , and  $-10R_j \leq y, z \leq 10R_j$ , where  $R_j$  is the initial jet radius (which is also the code distance unit). The Cartesian coordinate system has its origin at the center of the box and the jet flows through the x-axis. The jet is continuously injected into the bottom of the box [at  $\mathbf{r} = (-30R_j, 0, 0)$ ]. Inside the box, the SPH particles are initially distributed on a cubic lattice. An outflow boundary condition is assumed for the boundaries of the box. The particles are smoothed out by a spherically symmetric kernel function of width  $h$ , and the initial values of  $h$  were chosen to be  $0.4R_j$  and  $0.2R_j$  for the ambient and jet particles, respectively, so that we have up to 400,000 SPH particles at the beginning of the calculation.

We adopt a sinusoidal velocity profile for the pulsing jet at the inlet:

$$v_o(t) = v_j[1 + A \cdot \sin(\frac{2\pi}{P}t)] \quad (2)$$

where  $v_j$  is the mean jet speed, and  $A$  and  $P$  are the amplitude and the period of the velocity oscillations, respectively. We have chosen two periods for our models, namely,  $P = 1$  and  $P \sim 0.5$  (in code units, where the unit of time in our calculations corresponds to the ratio  $t_d = R_j/c_a \sim 38$  years, for the set of parameters adopted here; see below). The number density in the flow at injection,  $\rho_j$ , is assumed to be constant, which implies that  $\dot{m} = \rho_j \cdot v_j(t)$  is not constant. These models are equivalent to the constant density injection models recently investigated analytically by Cantó et al. (2000).

We consider different initial magnetic field profiles. As in papers I, II and III, we adopt 1) an initial constant longitudinal magnetic field [ $\mathbf{B} = (B_0, 0, 0)$ ], both inside and outside the jet beam; and 2) a helical force-free field that also extends to the ambient medium which is described by the following equations (see Todo et al. 1993, and Fig. 1 of Paper II):

$$B_r = 0 \quad (3a)$$

$$B_\phi(r) = B_0 \left[ \frac{0.5Cdr^2}{(r + 0.5d)^3} \right]^{1/2} \quad (3b)$$

$$B_x(r) = B_0 \left[ 1 - \frac{Cr^2(r + d)}{(r + 0.5d)^3} \right]^{1/2} \quad (3c)$$

where  $r = \sqrt{y^2 + z^2}$  is the radial distance from the jet axis and the  $C$  and  $d$  constants are given by 0.99 and  $3R_j$ , respectively. In these equations,  $B_0$  is the maximum strength of the magnetic field and corresponds to the magnitude of the longitudinal component at the jet axis. For these configurations, the jet is assumed to have an initially constant gas pressure ( $p_j$ ) which is in equilibrium with the ambient gas pressure ( $\kappa = p_j/p_a = 1$ ).

The third adopted configuration is a purely toroidal magnetic field ( $B_\phi$ ) whose functional form is given by Fig. 1 and equation (5) of SH. Only the jet beam is initially magnetized in this case, and for the sake of clarification, we here rewrite SH's equation for  $B_\phi$ :

$$B_\phi(r) = \begin{cases} B_{\phi,m} \frac{r}{r_m} & 0 \leq r \leq r_m \\ B_{\phi,m} \frac{R_j - r}{R_j - r_m} & r_m \leq r \leq R_j \\ 0 & R_j < r \end{cases} \quad (4)$$

where  $r_m$  is a free parameters that determines the position where the magnetic field has its maximum intensity (we adopt here  $r_m \sim 0.9R_j$ ). Lind et al. (1989), Frank et al. (1997, 1998) and OR, have also used a similar toroidal profile. In this case, in order to ensure initial magnetostatic equilibrium, the jet gas pressure has a radial profile with a maximum at the jet axis  $p_j(0) \approx 2.07p_j(R_j)$ , and  $p_j(R_j) = p_a$  at the jet surface (see Figure 1 of SH).

The physical conditions at the jet inlet are parameterized by the following nondimensional numbers:  $\eta = n_j/n_a$ , the jet-to-ambient number density ratio;  $M_a = v_j/c_a$ , the average jet velocity to the ambient sound velocity ratio;  $\kappa = p_j/p_a$ , the jet-to-ambient thermal pressure ratio;  $\beta = 8\pi p_{th}/B^2$ , the thermal-to-magnetic pressure ratio. As we will also study jets in which the magnetic field is a function of the radial distance  $r$ , we can define a mean magnetossonic Mach number as being  $\langle M_{ms} \rangle = \int M(r)rdr / \int rdr$ , where  $M(r) = v_j/[c + v_A(r)]^{1/2}$  (e.g., SH), and  $v_A$  is the Alfven speed (see Table 1).

## 2.2. Magnetic Field Reversals

We should make some remarks on the late evolution of radiatively cooling magnetized jets containing longitudinal components. In Paper II, we have shown examples with the development of magnetic field reversals at the contact discontinuity between the jet and the cocoon with intensities up to 5 times their initial magnitude. As stressed in Paper II, field reversals of the longitudinal component occur in both purely longitudinal and helical magnetic field configurations, because the field lines are amplified by compression in the nonparallel shocks at the jet head, and are enforced to flow backward with the shocked plasma into the cocoon. In this process, the lines are reoriented and sometimes have their polarization reversed. The increasing strength of the reversed fields due to shear at the contact discontinuity could lead to the development of strong pinching regions in the late evolution of some flows that ultimately could cause jet disruption (see Fig. 16 of Paper II). Although shear and compression are expected to enhance  $\mathbf{B}$ , we have speculated in Paper II that very large amplifications of the reversed fields could be due either to inappropriate computation of field dissipation under our ideal-MHD treatment, or to numerical effects.

In order to test the first of these hypotheses, we have carried out numerical simulations under a non-ideal MHD approximation, using different values for the magnetic resistivity. Although a complete analysis of the results of these tests is out of the scope of the present paper, we can briefly comment their consequences to the present study. We have found that a non-null magnetic resistivity ( $\eta_M$ ) is able only to postpone the effects above at the jet/cocoon interface [tested for a broad range of values of  $\eta_M$  from the thermal value (which is  $\propto T^{-3/2}$ ; e.g., Spitzer 1956), to values many orders of magnitude higher].

These findings have suggested a possible numerical origin for the large field reversals found in the earlier calculations. Performing then, extremely high resolution SPMHD calculations, using up to  $10^6$  SPH particles and the same conditions as in Fig. 16 of Paper II, we could eliminate

the anomalous amplification of the reversed fields and the disruption of the beam, while the other features that appeared in the original lower resolution model were left unaffected both at the head and along the jet beam. The anomalous amplification of the fields was, therefore, being caused essentially by a poorer interpolation of the physical quantities at the contact discontinuity in the lower resolution models. As the extremely high resolution simulations expend too much computational time and memory without significantly improving the overall results, we have instead, presently introduced in our code a numerical *cleaning switch*, in order to eliminate the anomalous magnetic fields at the jet/cocoon interface, in calculations involving longitudinal magnetic fields. The cleaning switch allows the reduction of the strength of the magnetic fields at the jet/cocoon interface by an amount of about 1% whenever the reversed  $\mathbf{B}$ -field component exceeds a critical value that was chosen to be twice as large as the maximum value of the field in the head (where  $\mathbf{B}$  regularly reaches its maximum physical amplification behind the strong shocks). In average, less than 0.01% of the SPH particles have their magnetic field vectors modified by such switch, which is applied in only less than 1% of the total number of time steps. The tests have indicated that the system is stabilized against anomalous field amplification and jet disruption with the employment of this switch and, more important, the results of these models are comparable with those obtained in the extremely high resolution calculations.

### 3. The simulations

In this section, we present the results of our 3-D, MHD numerical simulations of radiatively cooling pulsed jets. The physical parameters of our models are summarized in Table 1. All the parameters are suitable for the study of protostellar jets, namely:  $R_j \sim 2.5 \times 10^{15}$  cm,  $M_a \approx 15$ , or  $v_j \approx 15 \times c_a \approx 250$  km s $^{-1}$ ;  $\eta = n_j/n_a = 5$ ;  $n_a = 200$  cm $^{-3}$  (e.g., Reipurth & Raga 1999). We have adopted in most of the simulated models, a maximum intensity for the magnetic field as given by the equipartition condition:  $\beta = 8\pi p_{th}/B^2 \simeq 1$ , which implies a maximum intensity of  $\approx 80$   $\mu$ G. In some cases, we have adopted  $\beta \approx 0.1$ , which implies a maximum  $B \approx 260 \mu$ G. The MHD models are all compared with the purely HD counterparts for which  $\beta = \infty$ .

#### 3.1. Effects of different magnetic field configurations

Figure 1 displays the midplane density contours (left) and the velocity field distributions (right) for four supermagnetosonic, radiatively cooling, pulsed jets with initial  $\beta \simeq 1$  after they have propagated over a distance  $\approx 35R_j$ , at  $t/t_d \simeq 3$ . The top jet is purely hydrodynamical (model HD1 in Table 1); the second jet (from top to bottom) has an initial constant longitudinal magnetic field (model ML1 in Table 1); the third jet has an initial helical magnetic field (MH1), and the bottom jet has an initial toroidal magnetic field (MT1). The early evolution of these jets was discussed in Paper III.

At the time depicted in Figure 1, the leading working surface at the jet head is followed by 4 other features. Each new pulse at the left forms within a time of about  $0.2t_d$ . This value agrees with that predicted by linear analysis of pulsed jets with sinusoidal injection profile (Raga et al. 1990; Raga & Cantó 1998). Like the leading working surface, each internal feature consists of a double-shock structure, an upstream reverse shock that decelerates the high velocity material entering the pulse, and a downstream forward shock sweeping up the low velocity material ahead of the pulse. Each of these internal working surfaces (IWS) or knots, widens and broadens as it propagates downstream and squeezes shocked material sideways into the cocoon. Since  $P \simeq 0.5t_d$ , there has been time for 5 IWSs to form through steepening of the input sinusoidal profile. The first of these IWSs has already overtaken the leading working surface and merged with it. With the impact, it is disrupted and its debris are partially deposited into the cocoon thus providing a complex fragmented structure at the head as required by the observations (as is the case, for example, of HH 1/2; e.g., Jeff Hester, Stapelfeldt & Scowen 1998).

As previously reported (see Paper III, and de Gouveia Dal Pino & Cerqueira 2001), Fig. 1 suggests that the overall morphology of the 3-D pulsed jet is not very much affected by the presence of the different magnetic field configurations in equipartition with the gas. Compared with the purely HD jet (top panel), the introduction of the distinct **B**-profiles tends, instead, to alter essentially the detailed structure and the emission properties (see below) behind the shocks at both the head and internal knots, particularly when helical or toroidal fields are present. For example, like in steady-state calculations (see Paper I and II), while the MHD jet with initial longitudinal **B**-field (second panel) exhibits a fragmented dense shell at the head which is very similar to that in the pure HD jet, the MHD jets with helical and toroidal fields present a more collimated head structure due to the action of the tension forces associated with the toroidal component of the **B**-field (third and fourth panels; see also Paper III).

Several physical quantities along the beam axis of the jets of Fig. 1 with initial longitudinal (ML1), and toroidal (MT1) fields are displayed in Figures 2a and 2b, respectively. As previously reported for the helical jet of Fig. 1 (see Paper III), we see (top panel, Figs. 2a and b) that the initial sinusoidal velocity profile impressed on the flow at the inlet steepens into the familiar sawtooth pattern (e.g., Raga & Kofman 1992) as the faster material catches up with the slower, upstream gas in each pulse. The longitudinal field ( $B_x$ ) depicted in Fig. 2a (bottom panel) is anti-correlated with the sharp density peaks within each IWS (second panel), where it attains a minimum value, and amplifies between them. On the other hand, the toroidal magnetic field ( $B_\phi$ ) depicted in Fig. 2b (bottom panel) sharpens within the knots and rarefies between them. These results are in agreement with those found in Paper III for the jet with initial helical geometry (see also Gardiner & Frank 2000).

Figures 3a and b depict the same HD jet (HD1, Fig. 3a), and the MHD jet with helical field (MH1, Fig. 3b) of Figure 1, but both more evolved, i.e., after they have propagated over a distance  $\simeq 60R_j$  ( $t/t_d \simeq 5$ ). We note in both the development of a gentle wandering or wiggling along the jet axis which is much more pronounced in the MHD jet. This is excited by the kink mode

of the Kelvin-Helmholtz (K-H) instability. From the simulation we obtain a wavelength for the oscillation  $\lambda \approx 16R_j$ , which is in good agreement with that derived for the resonant wavelength of the fundamental kink mode of the K-H instability from the linear theory (e.g., Hardee, Clarke & Rosen 1997; Hardee & Stone 1997, Hardee & Norman 1988). This effect could provide a potential explanation for the wiggling features often detected in HH jets (e.g., HH 111 and HH 46/47; Reipurth et al. 1997; Heathcote et al. 1996).

Figure 4 displays the temporal evolution of the density of the leading working surface (left panels) and of the second IWS (right panels) for the four jets of Figure 1: HD1 (top), ML1 (second panel), MH1 (third panel), and MT1 (bottom). As reported in previous 3-D calculations of steady-state jets (e.g., GB93; Paper I and II), the density of the leading working surface in the HD1 model (top panel, left) shows a temporal variability that is caused by global thermal instabilities (GTI) of the strong, high velocity radiative shocks (e.g., Gaetz, Edgar & Chevalier 1988; GB93; Paper II). Consistently with GTI theory (GB93), the density of the shell oscillates with a period which is of the order of the cooling time behind the shocks,  $\tau \approx 2.2t_{cool}$  (where  $t_{cool} \approx 0.77t_d$  for the HD1 model as obtained from the initial conditions, or  $\tau \approx 1.7t_d$ ). The introduction of the different magnetic field configurations alters both the amplitude and the pattern of the density oscillations, particularly in the presence of the toroidal field (bottom panel, left), where the delay in the oscillations can be interpreted as being due to the increase of the cooling distance behind the shock caused by the amplification of the toroidal component (e.g., Paper I and II). In the right panels of Figure 4, we see that the density of the IWSs also varies in time due to the global thermal instabilities, with a period  $\tau \approx 2t_d$ , and the oscillation pattern is also more affected in the presence of the toroidal field. In all cases, the oscillation peak decays with time as a consequence of both the sideways deposition of the shocked material into the cocoon and the damping of the GTI (GB93). [This density decay is less obvious in the leading working surface (left panels) because of the continuous accumulation of material in the head due to the merging of the IWSs.]

### 3.2. Effects of different pulsation periods

Figure 5 displays the midplane density contours (left) and the velocity field distributions (right) for four pulsed jets after they have propagated over a distance  $\approx 40R_j$ , at  $t/t_d = 3$ . The top jet is purely hydrodynamic (model HD2 in Table 1); the second (from top to bottom) has an initial constant longitudinal magnetic field (model ML2 in Table 1); the third has an initial helical magnetic field (MH2), and the bottom has an initial toroidal magnetic field (MT2). The initial conditions are the same as in Figure 1, but here the period of the injected sinusoidal velocity profile is twice as large ( $P = 1t_d$ ; see Table 1). The four jets in Figure 5 are slightly advanced in the computational domain when compared to those of Figure 1. This is a natural consequence of the larger period  $P$  of the injected velocity variability (see, e.g., de Gouveia Dal Pino 2001). As the average propagation velocity of the leading bow shock ( $v_{bs}$ ) is larger than that of the jets of Figure 1, a smaller amount of shocked material is accumulated behind the jet shock or Mach disk (whose

velocity is  $v_{Md} \simeq v_j - v_{bs}$ ). Therefore, the dense, cold shell that develops at the head due to the radiative cooling of the shocked material is thinner and less clumpy in the larger period jets.

As before, the overall jet morphology is not very much affected by the presence of the different magnetic field configurations, although the tension forces associated with the toroidal magnetic field component in the MHD jets with initially helical and toroidal geometries (third and fourth panels) increases the jet collimation with respect to the HD jet. Global thermal instabilities also occur in these models. As in Figure 4, Figure 6 shows that the leading working surface (top-left panel) and the IWSs (here represented by the second IWS, top-right panel) of the purely HD jet also undergo temporal density oscillations due to the GTI. Likewise, the shape and amplitude of the oscillations are affected by the introduction of the magnetic fields, particularly in the helical (third panels) and toroidal (fourth panels) cases.

The  $H\alpha$  intensity is known to be strong behind the knots of the observed HH jets. As discussed in Paper III, we can estimate the  $H\alpha$  emissivity ( $I_{H\alpha}$ ) behind the IWSs using the results of our simulations and the relation given by Raga & Cantó (1998),  $I_{H\alpha} \propto \rho_d v_s^{3.8}$ , where  $\rho_d$  is the downstream pre-shock density, and  $v_s$  is the shock speed. In Paper III (see also de Gouveia Dal Pino & Cerqueira 2001), we have found for the MHD jets of Fig. 1 that the  $H\alpha$  intensity behind the internal knots is amplified up to factors  $\sim 3$ , and 4 in the presence of the helical and toroidal field geometries, respectively, relative to the HD jet. The same evaluation performed for the jets of longer period of Figure 5, results in amplifications of the  $H\alpha$  intensity up to factors  $\sim 5$ , and 4, for the MHD jets with helical and toroidal fields, respectively, with respect to the  $H\alpha$  intensity in the HD jet, while for the MHD jet with initial longitudinal field, the  $H\alpha$  intensity is essentially the same as that of purely HD jet (see Figure 7).

### 3.3. Effects of different values of $\beta$

Figures 8 and 9 compare the early evolution of jets with different magnetic field strengths. Figure 8 displays two MHD jets with initial longitudinal magnetic field geometry: a  $\beta \simeq 1$  (ML1, middle panel), and a  $\beta \simeq 0.1$  jet (ML1b, bottom panel), which are compared with a HD jet (HD1,  $\beta = \infty$ , top panel). Likewise, Figure 9 displays two MHD jets with initial helical magnetic field geometries: a  $\beta \simeq 1$  (MH1, middle panel), and a  $\beta \simeq 0.1$  jet (MH1b, bottom), which are also compared with the HD jet (HD1, top panel). The remaining initial conditions in these figures are the same as in Figure 1. In the  $\beta \simeq 0.1$  models, the maximum initial magnetic field is  $\approx 260\mu\text{G}$ . For both magnetic field configurations, we note a considerable increase in the jet collimation with larger  $\mathbf{B}$ -field ( $\beta \simeq 0.1$ ) that is obviously caused by the action of the much stronger confining magnetic tension and pressure forces. Nonetheless, the global characteristics of the flow found in the HD, and  $\beta \simeq 1$  jets are retained in the  $\beta \simeq 0.1$  flows, particularly in the longitudinal field configuration (Figure 8). For example, the clumpy structures that develop at the head of the HD, and  $\beta \simeq 1$  jets as a consequence of the combined effects of the radiative cooling and Rayleigh-Taylor (R-T) instabilities (see Paper I and II), are still present in the  $\beta \simeq 0.1$  jet. This result is somewhat in

contradiction with the recent 2-D results of OR, who found that the growth of the R-T instability and, therefore, the development of fragmentation at the jet head, could be inhibited in the presence of longitudinal fields even at equipartition conditions. This could be due to the differences in the assumed input conditions of their model with respect to the present ones. We note, however, that according to R-T instability theory an enhancement of the instability is expected in the presence of longitudinal  $\mathbf{B}$ -field (e.g., Jun, Norman & Stone 1995).

Figure 10 depicts the velocity (top panel), density (second panel), toroidal (third panel), and longitudinal (bottom panel) magnetic field components along the beam axis <sup>3</sup> for the  $\beta \simeq 1$  jet of Figure 9 (model MH1b in Table 1), at  $t/t_d = 2$ . As before, the toroidal component of the magnetic field (third panel) sharpens within the knots and rarefies between them, while the longitudinal component (forth panel) is stronger between the knots. For the  $\beta \simeq 1$  models with helical and toroidal configurations (see Figs. 2b, and 9), we find that the toroidal magnetic field components behind the IWSs are amplified by factors of  $\lesssim 1.5$  with respect to the initial values. For the  $\beta \simeq 0.1$  model (Figs. 9 and 10), we find only a slightly larger amplification of  $B_\phi$  behind the IWSs  $\gtrsim 2$ .

As in previous 2-D calculations (SH, OR), we also find that the density behind the IWSs increases with increasing  $\mathbf{B}$ -strength (or decreasing  $\beta$ ), particularly in the presence of toroidal field components (Figure 9). This density enhancement is caused again by the confining tension forces associated with the toroidal fields which are larger for lower  $\beta$ -jets. However, in our 3-D calculations, we find smaller density amplifications. For example, for the  $\beta \simeq 1$ , and 0.1 jets with helical fields (Figure 9), we find a maximum density amplification behind the IWSs of a factor  $\lesssim 1.4$ , and  $\lesssim 4$ , respectively, relative to the HD jet, and for the  $\beta \simeq 1$  jet with purely toroidal field (Fig. 1), we find a maximum density amplification of a factor  $\lesssim 1.6$ . On the other hand, SH have found larger density amplifications by factors  $\sim 2 - 40$  in their 2-D simulations of jets with purely toroidal fields and  $\beta \simeq 1$ , and  $\simeq 0.25$ , respectively. In the 2-D simulations of OR, the magnetic compression of the IWSs by the toroidal fields is more drastic, causing the disruption of them, even in  $\beta \simeq 1$  jets. Although these differences can be partially attributed to differences in the input conditions of each model (see §1), the results above suggest that the effects of the magnetic fields upon the jet structure tend to be smoothed out in 3-D flows (see discussion below).

#### 4. Discussion and Conclusions

We have presented the results of fully three-dimensional (3-D), magnetohydrodynamic (MHD) numerical simulations of radiatively cooling, overdense, pulsed jets, using the Smoothed Particle Magnetohydrodynamic (SPMHD) technique. We have explored the role of magnetic fields in pulsed jets considering different magnetic field strengths ( $\beta \simeq 0.1, 1$ , and  $\infty$ ) and configurations: longitu-

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<sup>3</sup>The toroidal magnetic field component,  $B_\phi$ , has been actually taken close to the jet surface other than along the jet's central axis where its value is too small (see eq. 3b).

dinal (aligned with the jet axis), helical, and purely toroidal magnetic fields. Our results could be summarized as follows:

1 - Magnetic fields with intensities of a few  $10\mu\text{G}$ , which are of the order of those inferred for Herbig-Haro (HH) jets (e.g., Morse et al. 1992, 1994), and which are in close equipartition of energy density with the thermal gas are able to improve collimation of the jet with respect to a non-magnetic jet. Such collimation effect is, as expected, more pronounced for magnetized jets with initial helical and toroidal fields.

2 - Although of relevance for the formation of structures in the shell at the jet head, the combined effect of non-uniform cooling and the Rayleigh-Taylor (R-T) instability (e.g., GB93; Paper I), which is particularly evident in the case of steady-state jets (see Paper I), becomes less important in pulsed jets. In these cases, the continuous impact of the internal working surfaces (IWSs) with the jet head can also cause shell fragmentation. While the presence of toroidal magnetic field components can modify the morphology close to the jet head inhibiting its R-T fragmentation in the early evolution of the jet, as previously reported (paper I, III), the later impact of the pulsed-induced internal knots causes the appearance of a clumpy, complex morphology at the jet head (as required by the observations of Herbig-Haro jets) even in MHD models with helical or toroidal configurations.

3 - The overall morphology of the jet does not seem to be very much affected by the presence of the different magnetic field configurations, or different magnetic field strengths (see also paper III). Nonetheless, the role of the magnetic fields in the determination of the detailed structure and the emission properties behind the shocks in the head and the IWSs cannot be ruled out. The 3-D simulations show that the density of the shell and IWSs undergoes oscillations with time (with periods of the order of the cooling time behind the shocks) that are caused by global thermal instabilities of the strong radiative shocks. The amplitude and shape of these oscillations are affected by the introduction of magnetic fields, particularly by toroidal configurations.

4 - The 3-D calculations also indicate that the the  $\text{H}\alpha$  emissivity can be enhanced by factors up to  $\sim 4$ , and 5 in the MHD models with toroidal field components, and  $\beta \simeq 1$ , and 0.1, respectively, when compared with the non-magnetic calculation. Using the same data from Hartigan, Raymond & Hartman (1987; see also Hartigan, Morse & Raymond 1994) that have allowed to obtain  $I_{\text{H}\alpha} \propto \rho_d v_s^{3.8}$  for the  $\text{H}\alpha$  intensity (Raga & Kofman 1992), we find that the intensity of the [S II] doublet is weakly dependent on the shock velocity ( $I_{[\text{S II}]} \propto v_s^{0.2}$ ). We can thus infer from our results and from these two relations, that the  $I_{[\text{S II}]} / I_{\text{H}\alpha}$  ratio (which is frequently used to determine the excitation level of the HH objects) can decrease by factors up to  $\sim 4 - 5$  in the MHD models with helical or toroidal fields (for  $\beta \simeq 1 - 0.1$ , respectively), with respect to the HD jet. Thus if the HH jets are really imbedded in magnetic fields with strengths not far from the equipartition with the gas, the evaluations of the physical quantities (such as, ionization fraction, mass loss ratio, etc.), which are presently made without taking into account the magnetic fields, could be significantly modified. We note, however, that the estimates above of  $I_{[\text{S II}]} / I_{\text{H}\alpha}$  are very crude and preliminary. Future

calculations coupling the (magneto)hydrodynamic calculation and the evolution of the chemical species will allow to obtain more precise estimates of the intensity ratios above. Also, it’s interesting to note that in the regime of IWS shock velocities, and preshock densities and  $\mathbf{B}$ -fields examined here ( $v_s \simeq 30 \text{ km s}^{-1}$ ,  $B < 300\mu\text{G}$ ,  $n_j \simeq 1000 \text{ cm}^{-3}$ ), previous “planar” shock calculations by Hartigan, Morse & Raymond (1994) have indicated that the  $I_{[\text{S II}]} / I_{\text{H}\alpha}$  ratio is less sensitive to the magnetic field intensity than in our calculations.

5 - The differences that arise in 2-D simulations with different magnetic field geometry and strengths seem to diminish in the 3-D flows. Like in 2-D models, the density behind the IWSs also increases with increasing  $\mathbf{B}$ -strength (or decreasing  $\beta$ ) in the 3-D calculations, but at a lower rate. For jets with initial helical field and  $\beta = 1$ , and 0.1, we have found maximum density amplifications behind the IWSs of factors  $\lesssim 1.4$ , and  $\lesssim 4$  respectively, and for a  $\beta \simeq 1$  jet with pure toroidal field, we have found a maximum density amplification of a factor  $\lesssim 1.6$ , relative to the HD jet. These amplification factors are substantially smaller than those found in the 2-D models (SH; OR).

6- The amplification of the toroidal field components behind the IWSs is also smaller in 3-D flows. We have found field amplification factors of  $\lesssim 1.5$  for MHD models with  $\beta \simeq 1$  and helical or toroidal configurations, and amplification factors  $\gtrsim 2$  in the  $\beta \simeq 0.1$  model with helical configuration, with respect to their initial magnetic fields. These values are much smaller than those detected, for instance, by SH, who found field amplifications up to  $\approx 7$  and 60 in their  $\beta \simeq 1$ , and 0.25 models, respectively.

7 - We have detected no signatures for the formation of nose cones. Such elongated structures are often found to develop at the head of magnetized jets in 2-D simulations involving toroidal magnetic fields. They are formed as the shocked gas is confined between the Mach disk and the bow shock by the tension forces associated with the toroidal field (e.g., Lind et al. 1989; Köss, Müller & Hillebrandt 1990; OR; SH). Although our 3-D calculations are probably limited by lower numerical resolution, the typical extensions of the nose cones in the 2-D calculations are few times the jet radius (SH; OR), so that they should be apparent even in a low resolution calculation. Our results thus indicate that nose cones are unstable in 3-D. Perhaps more convincing is the fact that the jet shock (or Mach disk) is found to be very close to the bow shock in most of the HH jets, as is the case of the HH 47A which, following the interpretation of Morse et al. (1994), may have a transverse magnetic field of the order of  $\approx 150\mu\text{G}$ .

8 - We speculate that the lower density amplification and the absence of nose cones in the 3-D calculations could be, in part, explained by the fact that, as the magnetic forces are intrinsically three-dimensional, they cause part of the material to be deflected in a third direction, therefore, smoothing out the strong focusing of the shocked material that is otherwise detected in the 2-D calculations (see also de Gouveia Dal Pino & Cerqueira 2001).

9 - Our simulations involving helical magnetic field geometry reveal the development of a wiggling structure at the late evolution that is due to the excitation of the kink mode of the Kelvin-Helmholtz instability. This could potentially explain the gentle wandering often displayed

by HH jets (e.g., HH 111, e.g., Reipurth et al. 1997; HH 46/47, e.g., Heathcote et al. 1996). The appearance of this mode whose development is not possible in 2-D axisymmetric models also reinforces the importance of fully 3-D calculations.

We are thankful to Alex Raga for his advice and suggestions. A.H.C. would like to thank the Brazilian agency FAPESP, that have supported this work under a Ph.D. Fellowship program (process 96/08182-3). E.M.G.D.P. acknowledges the Brazilian agencies FAPESP and CNPq for partial support. The simulations were performed on a cluster of Linux-based PC's, whose purchase was made possible by FAPESP. Also, partial support from the PRONEX/FINEP (41.96.0908.00) project is acknowledged.

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Fig. 1.— Midplane density contours (left) and the velocity field distributions (right) for four supermagnetosonic, radiatively cooling, pulsed jets after they have propagated over a distance  $\approx 35R_j$ , at  $t/t_d = 3$  (where  $t_d \equiv R_j/c_a \approx 38$  years). The top jet is purely hydrodynamic (model HD1 in Table 1); the second jet (from top to bottom) has an initial constant longitudinal magnetic field configuration (model ML1 in Table 1); the third jet has an initial helical magnetic field configuration (ML1), and the bottom jet has an initial toroidal magnetic field configuration (MT1). The initial parameters are the same for all the models:  $M_a = 15$ ,  $v_0 = v_j[1 + A \cdot \sin(2\pi t/P)]$ , with  $v_j \simeq 250 \text{ km s}^{-1}$ ,  $A = 0.25$  and  $P = 0.54t_d (\approx 20 \text{ years})$ ,  $\eta = 5$ ,  $\beta = \infty$  for the HD model, and  $\beta \simeq 1$  for MHD models. The maximum density in each model (from top to below) is:  $n/n_a \approx 156, 173, 140$  and  $183$  ( $1 n_a = 200 \text{ cm}^{-3}$ ). The  $x$  and  $z$  coordinates are in units of  $R_j$ . The jet is injected into the computational domain at  $x = -30R_j$ .

Fig. 2.— From top to bottom: velocity, density and magnetic field component profiles along the beam axis for the jets of Fig. 1 with: *a*) initial longitudinal magnetic field (ML1), and *b*) initial toroidal field (MT1). Time displayed  $t/t_d \simeq 3$ . The magnetic field intensities are displayed in code units (one code unit is  $\approx 21 \mu\text{G}$ ). Note that the coordinate along the jet axis ( $d$ ) has been shifted from  $-30R_j$  to 0. (The noise we see near the head results from the approximation of the first internal knot to the head.)

Fig. 3.— Midplane density contours (top panel) and the velocity field distributions (second panel) for: *a*) the HD1 model (top panel of Fig. 1), and *b*) the MH1 model (third panel of Figure 1), at  $t/t_d = 5$ . The bottom panel of Fig. 2b shows the magnetic field distribution.

Fig. 4.— Temporal evolution of the density in the leading working surface (left) and for the second internal working surfaces (right) for the four jets of Figure 1: HD1 (top), ML1 (second panel; from top to bottom), MH1 (third panel) and MT1 (bottom).

Fig. 5.— The same as in Figure 1 but with  $P = 1t_d$ .

Fig. 6.— The same as in Figure 4, but for the jets of Figure 5.

Fig. 7.— Ratio between the  $\text{H}\alpha$  intensity along the jet axis evaluated within the IWSs for the different magnetized jets of Figure 5 and the  $\text{H}\alpha$  intensity of the purely hydrodynamical jet: MHD model with longitudinal field (MHD L, squares), MHD model with helical field (MHD H, bullets) and MHD model with toroidal field (MHD T, stars). The  $d$  coordinate along the jet axis has been shifted from  $-30R_j$  to 0 (with respect to  $x$  the coordinate of Figure 5). The symbols mark the positions of the two prominent IWSs of the jets of Fig. 5.

Fig. 8.— Midplane density contours (left) and the velocity field distributions (right) for three supermagnetosonic, radiatively cooling, pulsed jets after they have propagated over a distance  $\approx 20R_j$ , at  $t/t_d \simeq 1.75$ . The top jet is purely hydrodynamical (model HD1 in Table 1); the middle and bottom jets have an initial constant longitudinal magnetic field configuration, with  $\beta \simeq 1$ , and  $\beta \simeq 0.1$ , respectively (models ML1, and ML1b of Table 1). The  $\beta = \infty$ , and 1 models have been

displayed also in Figure 1 of Paper III).

Fig. 9.— Midiplane density contours (left) and the velocity field distributions (right) for three supermagnetosonic, radiatively cooling, pulsed jets after they have propagated over a distance  $\approx 20R_j$ , at  $t/t_d \simeq 1.75$ . The top jet is purely hydrodynamic (model HD1 in Table 1); the middle and bottom jets have an initial helical magnetic field configuration, with  $\beta \simeq 1$ , and  $\beta \simeq 0.1$ , respectively (models MH1 and MH1b in Table 1; the  $\beta = \infty$ , and 1 models have been displayed also in Figure 1 of Paper III).

Fig. 10.— From top to bottom: velocity, density, toroidal magnetic field component, and longitudinal magnetic field profiles along the beam axis for the bottom-jet of Fig. 9 with initial helical magnetic field. Time displayed  $t/t_d \simeq 2$ . The magnetic field intensities are displayed in code units (one code unit is  $\approx 21 \mu\text{G}$ ). The coordinate along the jet axis ( $d$ ) has been shifted from  $-30R_j$  to 0.

Table 1. Physical parameters of jets.

Model	$M_a^a$	$M_{ms_j}$	$M_{ms_a}$	$\eta$	$\beta$	P
HD1	15	33.5	15	5	$\infty$	0.5
ML1	15	22.6	10.1	5	1	0.5
ML1b	15	9.7	4.2	5	0.1	0.5
MH1	15	23.2	10.4	5	1	0.5
MH1b	15	9.7	4.3	5	0.1	0.5
MT1	15	26	11.6	5	1	0.5
HD2	15	33.5	15	5	$\infty$	1
ML2	15	22.6	10.1	5	1	1
MH2	15	23.2	10.4	5	1	1
MT2	15	26	11.6	5	1	1

<sup>a</sup>This Mach number is evaluated for the average jet speed  $v_j = 250 \text{ km s}^{-1}$  (see equation 2).

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